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Penguins leaving the pole: bound-state effects in $B \to K^*\gamma$

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Applying perturbative QCD methods recently seen to give a good description of the two body hadronic decays of the B meson, we address the question of bound-state effects on the decay $B \to K^*\gamma$. Consistent with most analyses, we demonstrate that gluonic penguins, with photonic bremsstrahlung off a quark, change the decay rate by only a few percent. However, explicit off-shell b quark effects normally discarded are found to be large in amplitude, although in the standard model accidents of phase minimize the effect on the rate. Using an asymptotic distribution amplitude for the K^* and just the standard model, we can obtain a branching ratio of a few $\times 10^{-5}$, consistent with the observed rate.

1 Introduction

This note reports on potential bound state modifications to the analysis of $B \to K^*\gamma$, which is usually given solely in terms the on-shell subprocess $b \to s\gamma$.

The flavor changing neutral currents involved in the decays of the B into $K^*\gamma$ do not exist to leading order in the standard model, but can occur in second order in the Weak interaction via emission and reabsorbtion of W bosons [1]. These loop diagrams are often called "penguins," and their magnitude can be greatly modified by strong interaction effects [2, 3, 4].

There is recent further interest in these decays because additional penguin-like contributions could come from particles not in the minimal standard model [5, 6]. Contributions to B into $K^*\gamma$ decay from loops of non-standard-model particles (such as loops of supersymmetric particles that might be called "penguinos") would be a signal of their existence. To take advantage of this possibility, more precise study of the decay in the standard model needs to be undertaken.

The subprocess $b \to s\gamma$, taken as a free decay, is usually treated as the only flavor changing contribution leading to $B \to K^*\gamma$ [7]. However, bound state effects could seriously modify results coming from this assumption. Bound state effects include modifications due to the quarks being off shell in $b \to s\gamma$, contributions involving gluonic penguins or double (photon plus gluon) penguins, and contributions from annihilation diagrams. The latter involve no neutral flavor changing currents at all.

We use shall perturbative QCD (pQCD) in our analysis (see also [8]), a methodology we have previously applied [9, 10] to hadronic decays and semileptonic form factors of the B, inspired by Ref. [11]. Examples of the Feynman diagrams we calculate are given in Fig. 1. We require as input the effective vertices that result from the penguin diagram analyses [3, 4]. Thereafter our calculations are quite explicit and are detailed below.

A preview of our results is as follows: Diagrams involving the subprocess $b \to s\gamma$ do dominate, and keeping just contributions from the most commonly cited effective vertex gives close to the correct answer. However, some luck underlies the last statement. Since the internal heavy quark propagator can go off shell, there is an additional, independent, effective vertex [4] that can contribute. It does so with an amplitude whose magnitude is about 2/3 that from the usual vertex, and only because of fortunate phase relations is the magnitude of the sum nearly the same as for the usual vertex alone. Other diagrams are shown to be small, although they lead to a someday measurable few percent difference in the decay rates of the charged and neutral B into $K^*\gamma$ decay.

A diagram we are forced to omit for now is the double penguin, Fig 1e, where both a photon and a gluon come out of the loop, as the effective vertex it gives is not calculated.

2 Calculations

We now begin to describe our calculations in more detail. The penguins are represented by blobs in Fig. 1, which may be interpreted as an effective Hamiltonian, expanded as

$$H_{eff} = -4\frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \sum C_i(\mu) O_i(\mu). \tag{1}$$

The operators O_i are listed in references [3] and [4].

Consider the photonic penguin diagrams in Fig. 1a. If the incoming and outgoing quarks in $b \to s\gamma$ are on shell, there is only one relevant operator in H_{eff}. Refs. [4] contain two operators that can contribute, but using $i\not{D}q = m_q q$ they can be seen to be equivalent. We will write the operators so that only one contributes when the quarks are on-shell. The commonly used operator relevant to radiative b decay is then

$$O_7 = \frac{e}{16\pi^2}\; m_b \bar{s} \sigma^{\mu\nu} F_{\mu\nu} \frac{1}{2} \left(1 + \gamma_5\right) b. \tag{2}$$

The numbering is that of Ref. [3]; unfortunately the notations of [3] and [4] do not match. The other operator is

$$O_2' = \frac{e}{16\pi^2}\; \bar{s} \sigma^{\mu\nu} F_{\mu\nu} \frac{1}{2} \left(1 + \gamma_5\right) (i\not{D} - m_b) b, \tag{3}$$

where the numbering is from Ref. [4] and the prime reminds us that we have put the on shell part into the other operator (and that we changed the location of a factor $Q_d = -1/3$).

Let us here record the coefficients at the W mass scale,

$$C_2'(m_W) = \frac{x}{24(1-x)^4} \Big((1-x)(18x^2 - 11x - 1) + 2(3x-2)(5x-2)\ln x \Big) = -0.47, \tag{4}$$

$$C_7(m_W) = \frac{x}{24(1-x)^4} \Big((1-x)(8x^2 + 5x - 7) + 6x(3x-2)\ln x \Big) = -0.19, \tag{5}$$

where the numerical values are for $x \equiv m_t^2/m_W^2 = 4$, and also record how these operators evolve down to lower scale,

$$C_2'(\mu) = C_2'(m_W) - \frac{22}{81}\left(1 - \eta^{-2/\beta_0}\right) - \frac{11}{81}\left(\eta^{4/\beta_0} - 1\right), \tag{6}$$

$$C_7(\mu) = \eta^{-16/3\beta_0} \left[C_7(m_W) - \frac{58}{135} \left(\eta^{10/3\beta_0} - 1 \right) - \frac{29}{189} \left(\eta^{28/3\beta_0} - 1 \right) \right]. \tag{7}$$

Quantity η is $\alpha_s(\mu)/\alpha_s(m_W)$ and β_0 is $11 - (2/3)n_f$, and following standard practice we have neglected mixing with operators that give small effect. Both C_7 and C_2' increase in magnitude with decreasing scale μ.

We make the peaking approximation for ϕ_B, the distribution amplitude of the B meson, wherein

$$\phi_B(x) = \frac{f_B}{2\sqrt{3}} \delta(x_1 - \epsilon_B). \tag{8}$$

The decay constant of the B is f_B and x_1 is the light cone momentum fraction carried by the light quark. The mass of the B is given by $m_B = m_b + \bar{\Lambda}_B$ and $\epsilon_B = \bar{\Lambda}_B / m_B$. For the K^* distribution amplitude, we write

$$\phi_{K^*}(y) = \sqrt{3} f_{K^*} y_1 y_2 \tilde{\phi}_{K^*}(y). \tag{9}$$

The normalization is

$$\int_0^1 dy_1 \, \phi_{K^*}(y) = \frac{f_{K^*}}{2\sqrt{3}} \tag{10}$$

so that $\tilde{\phi}_{K^*}(y)$ is unity for the (super-)asymptotic distribution amplitude. We also make the approximation that $m_{K^*} = 0$. The second diagram of Fig. 1a is then zero.

The spin projection operators for the initial and final hadronic states are $\gamma_5(\not{p} - m_B)/\sqrt{2}$ and $\not{\xi}(\not{k} + (m_{K^*}))/\sqrt{2}$, respectively, with p, k, and q being the momenta of the B, K^*, and photon, and ξ the polarization vector of the K^*. Angular momentum conservation allows only transverse polarizations.

The result is

$$M_{\text{photonic penguin}} = -\frac{8G}{m_B \epsilon_B} \left(C_7(\mu) I - 2C_2'(\mu) I' \right) \left(p \cdot q \epsilon \cdot \xi + i \epsilon_{\mu\nu\alpha\beta} p^\mu q^\nu \epsilon^\alpha \xi^\beta \right), \tag{11}$$

where ϵ is the polarization of the photon and

$$G = C_F \frac{e\alpha_s}{4\pi} \frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* f_B f_{K^*} \tag{12}$$

with $C_F = 4/3$. Also

$$I = (1 - \epsilon_B) \int_0^1 dy_1 \, \tilde{\phi}_{K^*}(y) \, \frac{(1 - y_1)(1 + y_1 - 2\epsilon_B)}{y_1 - 2\epsilon_B - i0^+} \tag{13}$$

$$= (1 - \epsilon_B) \left(-\frac{1}{2} + (1 - 2\epsilon_B) \left[i\pi + \ln \frac{1 - 2\epsilon_B}{2\epsilon_B} \right] \right) \tag{14}$$

and

$$I' = \int_0^1 dy_1\,(1-y_1)\,\tilde{\phi}_{K^*}(y) = \frac{1}{2}, \tag{15}$$

where the integrated results are for the asymptotic distribution amplitude. The imaginary part come from an internal propagator going on shell. This is often taken as a signal that pQCD is inapplicable. However, if properties of the reaction overall dictate short distance propagation only, then pQCD can still be used [12]. This is the situation here, as discussed in [9] and [13].

Numerical results will be given after we have discussed what prove to be the subdominant contributions. Our understanding of why they are subdominant is helped by some order of magnitude estimates. The crucial expansion parameter is $1/\epsilon_B$ and its logarithms. Factors of ϵ_B come from the propagators, and can also be induced, depending on circumstances, by the factor $y_1 y_2$ in the K^* distribution amplitude.

In the photonic penguin diagrams, Fig. 1a, the gluon connects at the lower vertex to on-shell quarks and its propagator gives a factor proportional to $1/y_1\epsilon_B$, where y_1 and ϵ_B are the momentum fractions of the two light quarks. Thus appears one factor of $1/\epsilon_B$; the y_1 is canceled from the K^* distribution amplitude. The b quark propagator is involved in two subprocesses: scattering from the light quark by gluon exchange and decay into the s quark plus photon. Both are possible for an on-shell b quark, and the b quark does go on-shell in this diagram when $y_1 = 2\epsilon_B$. The b quark propagator thus contributes an imaginary pole term and a real principal value term (or just a real term, for the operator O_2') to the integral involving the K^* distribution amplitude, and one of them gives (roughly speaking) an $i\pi$ and the other gives a logarithm of $1/\epsilon_B$. Now we have accounted for the ϵ_B factors that appear in the photonic penguins,

$$M_{\text{photonic penguin}} \approx (\text{factors}) \times \frac{1}{\epsilon_B} \times \left(C_7(\mu) \;\times O\left(i\pi \text{ or } \ln\frac{1}{\epsilon_B}\right) + \; C_2'(\mu) \times O(1) \right). \tag{16}$$

The gluonic penguin graph with emission of the photon from the b or s quark (Fig. 1b) does not allow the quark propagator to be on shell. For example, in the diagram with photon emission from the s quark, the internal s quark decaying into an on-shell photon and and on-shell s quark must have a momentum squared larger than the mass squared of the s. A similar argument shows the b propagator is never on shell in the diagram with photon emission from the b quark. The gluon propagator is still spacelike and still gives a $1/\epsilon_B$, but compared to the previous case we lose the $\log(1/\epsilon_B)$ or $i\pi$ that came from quark propagator,

$$M_{\text{b or s bremss}} \approx (\text{similar factors}) \times \frac{1}{\epsilon_B} \times C_8(\mu) \times O\,(1)\,. \tag{17}$$

As we shall see, another significant reduction comes from the replacement of coefficient C_7 by its gluonic counterpart C_8.

For the spectator bremsstrahlung case, Fig. 1c, the quark propagator cannot go on-shell. However, the gluon propagator can. A factor $(1/\epsilon_B)$ that came from the gluon propagator is lost, and replaced by factors $i\pi$ or $\log(1/\epsilon_B)$ that come from integrating the K^* quarks's momentum fraction over the gluon pole, yielding

$$M_{\text{spectator bremss}} \approx (\text{similar factors}) \times (\epsilon_B)^0 \times C_8(\mu) \times O\left(i\pi \text{ or } \ln\frac{1}{\epsilon_B}\right). \tag{18}$$

Thus the gluonic penguin diagrams are suppressed by powers of ϵ_B or logs thereof, as well as by the ratio C_8/C_7.

For the actual calculations involving the gluonic penguin we kept just O_8 in the effective hamiltonian, where

$$O_8 = \frac{g}{16\pi^2}\; m_b \bar{s}\sigma^{\mu\nu} G_{\mu\nu} T_a \frac{1}{2}\left(1+\gamma_5\right) b. \tag{19}$$

The numbering is that of Ref. [3]. Other operators are possible when the gluon or quarks are off-shell. We have not explicitly calculated their contributions in this case (in part because their evolution has not been calculated), but have verified that the order of magnitude estimates are not upset, i.e., they are not leading in $1/\epsilon_B$.

The results are

$$M_{\text{b or s bremss}} = \frac{8e_d G}{m_B \epsilon_B} C_8(\mu) I' \left(p\cdot q \epsilon\cdot\xi + i\epsilon_{\mu\nu\alpha\beta} p^\mu q^\nu \epsilon^\alpha \xi^\beta\right), \tag{20}$$

where $e_d = -1/3$ and G and I' have the same meanings as before, and

$$M_{\text{spectator bremss}} = -\frac{4e_q G}{m_B} C_8(\mu) I_0(\epsilon_B) \left(p\cdot q \epsilon\cdot\xi + i\epsilon_{\mu\nu\alpha\beta} p^\mu q^\nu \epsilon^\alpha \xi^\beta\right), \tag{21}$$

where e_q is the quark charge for the spectator. We neglected some small terms in the numerator and let

$$I_0(\epsilon_B) = -i\pi + \ln\frac{1-\epsilon_B}{\epsilon_B}. \tag{22}$$

The coefficient C_8 is

$$C_8(m_W) = -\frac{x}{8(1-x)^4}\left((x-1)(x^2-5x-2) + 6x\ln x\right) = -0.094, \tag{23}$$

where the numerical value is again for $x = 4$, and it scales like,

$$C_8(\mu) = \eta^{-14/3\beta_0}\left[C_8(m_W) - \frac{11}{144}\left(\eta^{8/3\beta_0} - 1\right) + \frac{35}{234}\left(\eta^{26/3\beta_0} - 1\right)\right]. \tag{24}$$

As the renormalization scale decreases, the magnitude of C_8 decreases, actually passing through zero at $\mu \approx 6$ GeV.

Additionally, there are the annihilation graphs, Fig. 1d. These can contribute only to $B^{\pm}$ decay. To leading order in $(\epsilon_B)^{-1}$ the result comes from bremsstrahlung off the initial u quark and gives

$$M_{\rm ann} = \frac{2e_u}{m_B\epsilon_B}\frac{m_{K^*}}{m_B}\left[e\frac{G_F}{\sqrt{2}}V_{ub}V_{us}^* f_B f_{K^*}\right]\left(p\cdot q\epsilon\cdot\xi + i\epsilon_{\mu\nu\alpha\beta}p^\mu q^\nu \epsilon^\alpha \xi^\beta\right), \tag{25}$$

where we kept m_{K^*} when it appeared as an overall factor and where $e_u = 2/3$. The quantity in square brackets differs from the quantity G used earlier in lacking the strong interaction factors $C_F\alpha_s/4\pi$ and in having different CKM factors. While it is interesting that the decay proceeds at all without flavor changing neutral currents, the result turns out small. Not having the gluon exchange is a plus numerically, but the factors m_{K^*}/m_B and V_{ub} ensure the small result.

3 Numerical results

The numerical results should not depend on the renormalization scale. However, as we are most familiar with the wave functions or distribution amplitudes at a typical hadronic scale, say $\mu \approx 1$ GeV, we should evaluate the other quantities at the same scale. There are very big changes in the C_i from their values at the W mass scale.

For the sake of definiteness we shall use

$$\Lambda_{QCD} = 100 \text{ MeV},$$

$$\bar\Lambda_B = 500 \text{ MeV},$$

$$m_W = 81 \text{ GeV},$$

$$m_t = 2m_W,$$

$$V_{ts} = -0.045,$$

$$V_{tb} = 0.999,$$

$$V_{ub} = 0.0045,$$

$$V_{us} = 0.22,$$

$$\tau_B = 1.46 \text{ picoseconds},$$

$$f_B = 132 \text{ MeV},$$

and

$$f_{K^*} = 151 \text{ MeV}$$

Our convention has $f_\pi = 93$ MeV and the signs of the CKM parameters follow a "standard" advocated in [14]. The sign of $V_{tb}V_{ts}^*/V_{ub}V_{us}^*$ is what is crucial and it does not depend on conventions. We also use the asymptotic form for the K^* distribution amplitude and will mention results with another form later. We will express each contribution to the amplitude as

$$M_i = t_i \times \frac{1}{2}\left(\epsilon \cdot \xi + i(p \cdot q)^{-1}\epsilon_{\mu\nu\alpha\beta}p^\mu q^\nu \epsilon^\alpha \xi^\beta\right), \tag{26}$$

whereupon

$$\Gamma = \frac{1}{16\pi m_B}|t|^2 \tag{27}$$

for t being the sum of the t_i (and neglecting the K^* mass).

The scale we should use should be compatible with the scale that our wave functions and distribution amplitudes are determined at, and this is turn should be consistent with the scale of the four-momentum squared of the off-shell gluon. This suggests $\mu \approx 0.8$ GeV, which we shall use. We extrapolate the coefficients according to the formulas given earlier. Much of the change due to the extrapolation occurs as the scale changes from m_W down to order m_B, at least for the large terms C_7 and C_2'. For the running coupling we use $\alpha_s = 4\pi/\beta_0 \ln(\mu^2/\Lambda_{QCD}^2)$ with number of flavors appropriate to the scale.

From operators O_7 and O_2' we get

$$\begin{aligned} t_7 &= -0.95 - 3.56i \\ t_{2'} &= 2.40 \end{aligned} \tag{28}$$

with the contribution from photonic penguins being the sum of these two. The amplitudes are in units of 10^{-8} GeV. For others we get

$$\begin{aligned} t_{\text{b or s bremss}} &= 0.08, \\ t_{\text{spectator bremss}} &= 0.04 - 0.05i, \\ t_{ann} &= 0.06. \end{aligned} \tag{29}$$

The last two are given for the charged B. The photonic penguins dominate, although the other graphs contribute circa 10% corrections to the real parts of the decay amplitude.

The extra operator that we have considered in the photonic penguin calculation gives an amplitude with a magnitude that is about 2/3 as large as the amplitude from the operator normally considered. It is the luck of the phases that keeps its effect small. It increases the rate by not even 10% since roughly speaking what it does is just change the sign of the real part of the amplitude.

Specifically, we get a branching fraction

$$Br(B \to K^* \gamma) = 1.24 \times 10^{-5} \tag{30}$$

including just the photonic penguin terms. Keeping only the most usual operator O_7 would reduce the branching fraction to 1.13×10^{-5}. It seems inconsistent to include the smaller contributions since they may be smaller than the errors induced by our approximations upon the big terms. However, keeping all terms anyway gives 1.24×10^{-5} (unchanged from above) for the neutral B and 1.31×10^{-5} for the charged B. The relative size of the neutral and charged B decays should be about right and would be interesting to observe as more precise data becomes available.

It is possible that the distribution amplitude for the transversely polarized K^* is narrower than the asymptotic one. Chernyak, Zhitnitsky, and Zhitnitsky [15] suggested a distribution amplitude

$$\tilde{\phi}_{K^*}(y) = 5(y_1 y_2)^2, \tag{31}$$

albeit this was for the transverse ρ. If we use this distribution amplitude for the K^*, our calculated branching fractions are roughly halved.

The choice of Λ_{QCD} was made consistent with some of our own earlier work [9, 10], but could be varied (the earlier situation is much less sensitive to the value of this quantity than the present case will prove to be). If we let $\Lambda_{QCD} = 200$ MeV, leaving other parameter choices untouched, the branching ratio with the asymptotic distribution amplitude changes to

$$Br(B \to K^* \gamma) = 3.50 \times 10^{-5} \tag{32}$$

including just photonic penguins terms, with commensurate changes in results keeping all terms or just O_7. These values are in accord with present experimental data [16]. It is also clear that values for other parameters could be varied somewhat from values that we have used.

4 Conclusion

It seems with present knowledge, the actual $B \to K^*\gamma$ decay rate is sensitive to parameters of the bound state and to parameters governing the evolution of QCD. Still, a number of conclusions may be drawn.

Contributions from gluonic penguin and annihilation diagrams—which contribute to the physical decay but not to $b \to s\gamma$—have been calculated here. They change the decay rate by a few percent and so are not worrisome until the experiments are considerably more precise.

Also calculated here, and more significant, are effects due to decay quarks being off-shell. This brings into play another operator in the effective hamiltonian for $b \to s\gamma$, and this new operator produces an amplitude of noticeable magnitude. However, its phase is such that the effect on the decay is under 10%.

Regarding the future, there is a need to calculate the double penguin diagrams mentioned in the introduction and illustrated in Fig. 1e, including the QCD corrections to them. For now, we can at least note that these diagrams will not contribute an imaginary part to the amplitude, so that the imaginary part that is already there puts a lower bound on the total result. Also, one will wish to eventually study the totality of $B \to X_s\gamma$ since this is closer to the basic $b \to s\gamma$ than any individual exclusive channel. Still, the physical $B \to X_s\gamma$ always involves spectators and always has a contribution from operators that only contribute if the b is off-shell. We have not calculated for situations other than the lowest "X_s" and do not know what size corrections ensue overall, or even if the phase situation for the off-shell contributions persists.

None-the-less, the opportunity to test the flavor changing neutral currents induced in the standard model and to search for evidence of particles or phenomena beyond the standard model makes B decay into $K^*\gamma$ and into $X_s\gamma$ interesting, and makes calculations to determine precisely the standard model contributions to these decays a worthwhile and necessary pursuit.

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Figure caption

Fig. 1. Diagrams. The effective vertex, due for example to a W and t quark loop, is represented as an oval blob; (a) shows the photonic penguins, (b) shows gluonic penguins with bremsstrahlung from the b or s quark, (c) shows gluonic penguins with bremsstrahlung from the spectator quark, (d) shows two of four lowest order annihilation diagrams that could give charged $B \to K^*\gamma$, and (e) shows the "double penguin," once as a blob and once showing one example of a contribution to the blob.

Figure 1: